TeV leptogenesis in Z-prime models and its collider probe

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Abstract

We show that the $U(1)'$ models linked with the seesaw mechanism at TeV scale can lead to a successful baryogenesis through soft leptogenesis with a resonant behavior in the B parameter. Such a consideration constrains the Z' mass to be larger than $2-3$ TeV depending on the seesaw scale and the spharelon rate. Together with multi-TeV Z', large sneutrino-antisneutrino mixing and CP violating phenomena required by TeV leptogenesis could be searched for in future colliders by observing the distinct same-sign dilepton–dichargino as well as dislepton–diHiggs signatures.

The observed neutrino masses and mixing could be explained by the celebrated seesaw mechanism [1] at TeV scale, which requires the neutrino Yukawa coupling of the order of the electron Yukawa coupling. Such a seesaw mechanism may well be linked with an extended gauge symmetry spontaneously broken at low energy. A prototype would be the left-right symmetric model based on the gauge group $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ [2]. More economically, one can think of the gauge group $SU(2)_L \times U(1)_Y \times U(1)'$ which appears as a low-energy sector of the grand unification; $E_6 \rightarrow SU(5) \times U(1)_\chi \times U(1)_\psi$ where the $U(1)'$ is a certain linear combination of $U(1)_\chi$ and $U(1)_\psi$ [3]. In this case, three right-handed neutrinos N are usually required for anomaly cancellation [4]. Then, the TeV-scale seesaw mechanism can be directly tested by observing clean signatures of same-sign dileptons arising from the Majorana nature of the neutrinos [5], together with collider searches for the associated extra gauge boson, say, Z' of $U(1)'$.

Leptogenesis is an attractive feature of the seesaw mechanism, providing a nice way of generating the baryon asymmetry of the Universe [6]. In the usual thermal leptogenesis, the right-handed neutrino mass is required to be larger than about 10^8 GeV assuming the typical hierarchical mass pattern [7]. For a TeV-scale leptogenesis to work, one could however have highly degenerate right-handed neutrinos realizing a resonant enhancement [8]. In supersymmetric theories, new ways for leptogenesis arise due to lepton number violating (right-handed) sneutrino-antisneutrino ($\tilde{N}$-$\tilde{N}^\dagger$) mixing and CP violation in soft supersymmetry breaking sector [9]. This mechanism (so-called soft leptogenesis) can work within a single generation and requires unconventionally small B parameter of the lepton number violating soft mass. Such a situation may be realized in a radiative way consistently with TeV seesaw mechanism [10]. In gauge-mediated supersymmetry breaking theories, the small gravitino mass, $m_{3/2} = (10^{-2} - 10)$ eV, can be more naturally related to the similarly small B [11], which is indeed in the right range for our consideration. Recently, a number of attempts have been made to realize successful thermal leptogenesis at TeV scale [12, 13, 14], in anticipation of future detection.

In this paper, we will show that the supersymmetric seesaw mechanism associated with an extra gauge symmetry at TeV scale is a viable option for the generation of the cosmological baryon asymmetry, which leads to testable predictions in colliders. As mentioned, the existence of a new gauge interaction certainly makes the model highly testable as N and $\tilde{N}$

can be produced in colliders through, e.g., $e^+e^- \to NN, \tilde{N}\tilde{N}^\dagger$, which would not be the case only with the usually small neutrino Yukawa couplings, $h \sim 10^{-6}$. In favor of collider tests, one would like to have a larger gauge coupling and smaller masses for the new particles. However, such a situation may be in contradiction with the observed baryon asymmetry, $Y_B \sim 10^{-10}$. Indeed, the inverse process, i.e., the gauge annihilation, $\tilde{N}\tilde{N}^\dagger, NN \to$ light fermions, with the standard gauge coupling strength strongly suppresses the resulting lepton asymmetry. This puts a constraint on the new gauge boson mass, which depends sensitively on the spharelon dynamics below the electroweak phase transition [14, 15]. The existence of an extra gauge boson at mutli-TeV region will imply that the resonant behavior $B \sim \Gamma \sim 0.1$ eV and large CP violation have to occur for a successful leptogenesis. This is an interesting possibility for the collider search, as the model predicts testable signatures in the $\tilde{N}$-$\tilde{N}^\dagger$ oscillation and the associated CP violating phenomena, which may link cosmology with collider physics.

For the definite discussion of the seesaw mechanism with an extra gauge boson, let us take the Z' model with $U(1)_\chi$ which appears as $SO(10) \to SU(3)_c \times SU(2)_L \times U(1)_Y \times U(1)_\chi$, and assume the standard gauge coupling strength, $\alpha_\chi = \alpha' \approx 1/60$. In terms of the $SU(5)$ notation, the SM fermions and the singlet field N carry the $U(1)_\chi$ charges as follows:

$$\mathbf{10}\,(\frac{-1}{\sqrt{40}}), \quad \overline{\mathbf{5}}\,(\frac{3}{\sqrt{40}}), \quad \mathbf{1}\,(\frac{5}{\sqrt{40}}).$$

Furthermore, we will suppress the family indices of three right-handed (s)neutrinos, as the soft leptogenesis does not require mixing among different families, and then calculate the baryon asymmetry arising from a right-handed sneutrino decay with the typical neutrino Yukawa coupling corresponding to the atmospheric neutrino mass, $m_\nu = 0.05$ eV $\sim h^2 v^2/M$. The essential features of our results can be applied also to other classes of Z' models.

The superpotential of the seesaw model is

$$W = hLH_2N + \frac{1}{2}MNN \tag{1}$$

where L, H_2 and N denote the lepton, Higgs and right-handed neutrino superfields, respectively. The soft supersymmetry breaking terms are

$$V_{soft} = \left[Ah\tilde{L}H_2\tilde{N} + \frac{1}{2}BM\tilde{N}\tilde{N} + h.c.\right] + \tilde{m}^2\tilde{N}^\dagger\tilde{N} \tag{2}$$

where all the fields are understood as the scalar components of the superfields in Eq. (1), and $A, B, \tilde{m}$ are dimension-one soft parameters. To denote the corresponding fermion components, the notation of $L, \tilde{H}_2$ and N will be used. Without loss of generality, one can take M and B to be real and positive. Then, the right-handed sneutrino field is written as $\tilde{N} = (\tilde{N}_1 + i\tilde{N}_2)/\sqrt{2}$ in which two mass eigenstates $\tilde{N}_{1,2}$ have

$$M_{\tilde{N}_{1,2}}^2 = M^2 + \tilde{m}^2 \pm BM , \tag{3}$$

and thus the mass-squared difference $\Delta M_{\tilde{N}}^2 = 2BM$. For $B \ll M$, the mass difference becomes $\Delta M_{\tilde{N}} \approx BM/M_{\tilde{N}}$ where $M_{\tilde{N}}^2 \equiv M^2 + \tilde{m}^3$. The Yukawa couplings of $\tilde{N}_{1,2}$ can be read off from the Lagrangian

$$\mathcal{L} = h\tilde{N}L\tilde{H}_2 + h(A\tilde{N} + M\tilde{N}^\dagger)\tilde{L}H_2 + h.c. . \tag{4}$$

Let us now calculate the lepton and CP asymmetry ϵ_X arising from a particle X decaying to the final state F:

$$\epsilon_X \equiv \frac{\Gamma(X \to F) - \Gamma(X^\dagger \to F^\dagger)}{\Gamma(X \to F) + \Gamma(X^\dagger \to F^\dagger)}$$

where $F = LH_2$ or $\tilde{L}\tilde{H}_2$ in our case. Using the effective field theory approach of resumed propagators for unstable particles ($X = \tilde{N}_{1,2}$) [8], we find

$$\epsilon_{1,2} = \frac{2BMM_{\tilde{N}}(\Gamma_1 + \Gamma_2)\,\xi_{1,2}}{4B^2M^2 + M_{\tilde{N}}^2(\Gamma_{2,1}^2 - \Gamma_1\Gamma_2\xi_{1,2}^2)} \frac{M_{\tilde{N}}^2 + |A|^2 - M^2}{M_{\tilde{N}}^2 + |A|^2 + M^2}, \tag{5}$$

where

$$\xi_{1,2} = -\frac{|h|^2}{4\pi}\frac{\mathrm{Im}(A)M}{\Gamma_{1,2}M_{\tilde{N}}},$$

$$\text{and} \quad \Gamma_{1,2} = \frac{|h|^2}{8\pi}(1 + \frac{|A \pm M|^2}{M_{\tilde{N}}^2})M_{\tilde{N}} .$$

Note that, in Eq. (5), we have neglected the thermal effect breaking supersymmetry [9] which is suppressed compared to soft breaking effect by A and $\tilde{m}$ when they are of the same order of the right-handed neutrino mass M. As one can see, $\epsilon_{1,2}$ vanishes in the limit of $A \to 0$ and $M_{\tilde{N}} \to M$ without including the difference between the bosonic and fermionic thermal distributions. Note that one can have $\epsilon \sim 1$ for $B \sim \Gamma$ and the order one phase of A with $|A| \sim M, M_{\tilde{N}}$. Taking the approximation of $\Gamma \approx h^2M/4\pi \approx m_\nu M^2/4\pi v^2$, one obtains the resonance condition:

$$B \sim \Gamma \sim 0.1\mathrm{eV} \left(\frac{m_\nu}{0.05\ \mathrm{eV}}\right) \left(\frac{M}{\mathrm{TeV}}\right)^2 . \tag{6}$$

The Boltzmann equations governing the number densities of the sneutrino fields, Y_X, and lepton-antilepton asymmetry, Y_l, in unit of the entropy density are

$$\frac{dY_X}{dz} = -zK\left[\gamma_D(Y_X - Y_X^{eq}) + \gamma_A\frac{(Y_X^2 - Y_X^{eq\,2})}{Y_X^{eq}}\right]$$

$$\frac{dY_l}{dz} = 2zK\gamma_D\left[\epsilon(Y_X - Y_X^{eq}) - \frac{Y_X^{eq}}{2Y_l^{eq}}Y_l\right] \tag{7}$$

where $K \equiv \Gamma/H_1$ and $H_1 = 1.66\sqrt{g_*}M_{\tilde{N}}^2/m_{Pl}$ is the Hubble parameter at the temperature $T = M_{\tilde{N}}$. We take the Standard Model value of $g_* = 106.75$. In Eq. (7), $\gamma_D = K_1(z)/K_2(z)$ is the usual contribution from the decay due to the neutrino Yukawa coupling h, and γ_A accounts for the annihilation effect of the snuetrinos to the light Standard Model fermions mediated by the heavy Z' field: $\tilde{N}\tilde{N}^\dagger \to Z' \to f\bar{f}$. The gauge annihilation contribution is given by

$$\gamma_A = \frac{5}{\pi}\frac{\alpha_\chi^2 M_{\tilde{N}}}{K H_1}\int_1^\infty dt\,\frac{K_1(2zt)}{K_2(z)}\frac{t^3(t^2 - 1)^{3/2}}{(t^2 - \frac{1}{4}r^2)^2 + \frac{1}{16}u^2} \tag{8}$$

where $r \equiv M_{Z'}/M_{\tilde{N}}$ and $u \equiv r\Gamma_{Z'}/M_{\tilde{N}}$. For our numerical solution, we simplify the Boltzmann equation by taking $\Gamma = \Gamma_{1,2}$ and $\epsilon = \epsilon_{1,2}$. Taking the estimation of Γ in Eq. (6), we get $K \sim m_\nu/5\times10^{-4}\text{eV} \sim 100$ for $m_\nu = 0.05$ eV, independently of $M_{\tilde{N}}$. Note that the annihilation effect gets very strong with the factor of $M_{\tilde{N}}/H_1 \sim 10^{15}$ for the TeV-scale mass $M_{\tilde{N}}$. As one can expect from such large K and even larger gauge annihilation effect, the right-handed sneutrinos follow closely the thermal equilibrium distribution until very low temperature, and the lepton asymmetry freezes out at large z, that is, $T < 100$ GeV. Such behavior is shown in Fig. 1 where the evolution of the lepton asymmetry, $\log(Y_l/\epsilon)$, is plotted. One finds that Y_l becomes larger for a larger $M_{Z'}$ which suppresses more the annihilation effect, and finally approaches an asymptotic value for which the annihilation effect eventually drops out due to the Boltzmann suppression.

Let us note that the lepton asymmetry is still increasing during the electroweak phase transition $T \sim 100$ GeV which is well before the asymptotic value is reached. Therefore, it is important to include the spharelon effect after the phase transition to obtain the right amount of baryon asymmetry. To do this, let us define $\tilde{Y}_l = Y_L - Y_B$ where the $\tilde{Y}_l$ is the solution of the above leptogenesis equations (7) as given in Fig. 1. Then, it can be convoluted to the baryon asymmetry as

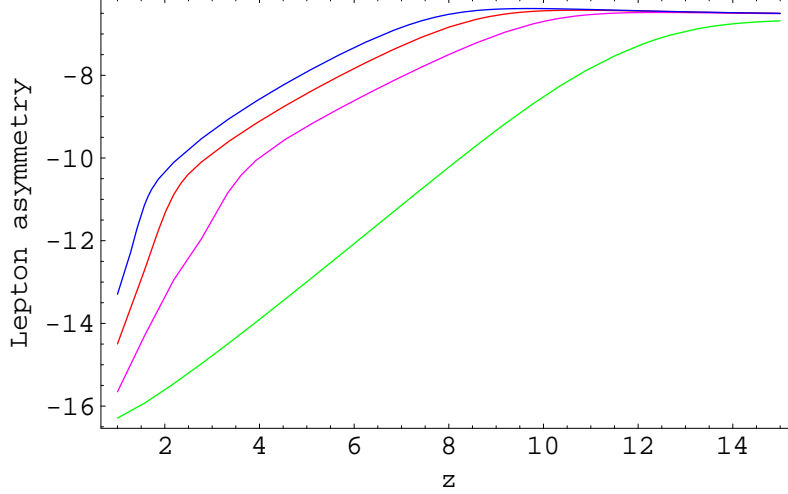$$Y_B(z) = -\int_0^z dy\, A_1(y)\tilde{Y}_l(y)\exp[-\int_y^z dx\, A_2(x)] \tag{9}$$

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FIG. 1: The lepton asymmetry in unit of ϵ is calculated as a function of $z = M_{\tilde{N}}/T$ with $M_{\tilde{N}} = 0.3$ TeV and $K = 500$. The curves are for $M_{Z'} = 1, 2, 3$ and 4 TeV from below.

where

$$A_1(z) = z\frac{\Gamma_{sp}}{H_1}\left(\frac{28}{51} + \frac{108}{561}\frac{v^2(T)}{T^2}\right)$$

$$A_2(z) = z\frac{\Gamma_{sp}}{H_1}\left(\frac{79}{51} + \frac{333}{561}\frac{v^2(T)}{T^2}\right).$$

Here Γ_{sp} is spharelon interaction rate and $v(T) = v_0(1-T^2/T_c^2)^{1/2}$ with $v_0 = 246$ GeV. Above the electroweak phase transition $T > T_c$, $v(T) = 0$ and Γ_{sp} can be taken to be infinitely large so that the standard result, $Y_B = -\frac{28}{51}Y_L = -\frac{28}{79}\tilde{Y}_l$, is recovered. Below T_c, several calculations for Γ_{sp} have been made within the validity range of $M_W(T) \ll T \ll M_W(T)/\alpha_w$ [15], which indicates that the spharelon interaction is still very active just below T_c and its freeze-out happens somewhat later [14]. One finds that Γ_{sp} is an extremely steep function of T and thus it is a fairly good approximation to calculate the final baryon asymmetry as follows:

$$Y_B \approx -\frac{A_1}{A_2}\tilde{Y}_l\Big|_{T_{sp}} \approx -\frac{1}{3}\tilde{Y}_l\Big|_{T_{sp}} \tag{10}$$

where T_{sp} is the spharelon freeze-out temperature. There is an uncertainty in determining T_{sp} which comes from the limited knowledge in calculating the spharelon rate. The most uncertain quantity, called κ, is known to lie in the range; $\kappa = (10^{-4} - 1)$ [15]. Adopting the parameters in Ref. [14], we find $T_{sp} = 80 - 90$ GeV which is in the region of $T > M_W(T)$ where the above spharelon calculation may well be extended.
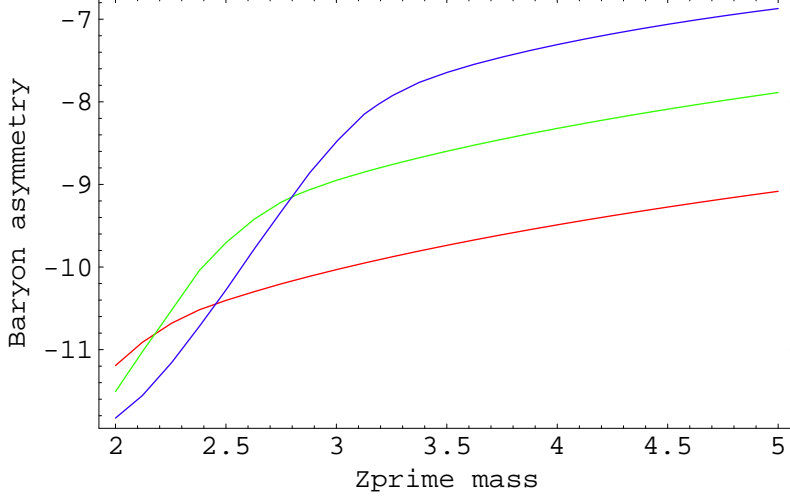
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FIG. 2: The baryon asymmetry in unit of ϵ is shown with respect to $M_{Z'}/\text{TeV}$ for $K = 100$. The curves are for $M_{\tilde{N}} = 0.3, 0.6$ and 0.9 TeV from below on the right-hand side.

Following the above prescription, the baryon asymmetry, $\log(Y_B/\epsilon)$, is calculated in Fig. 2 as a function of $M_{Z'}$ for various values of $M_{\tilde{N}}$, taking $T_{sp} = 90$ GeV. As one can see, Y_B depends strongly on the masses, $M_{Z'}$ and $M_{\tilde{N}}$, while we find it almost unaffected by the variation of K in the range: $K = 10 - 10^3$. Requiring $Y_B/\epsilon > 10^{-10}$, one obtains the lower bound: $M_{Z'} > (2.3 - 3)$ TeV for $M_{\tilde{N}} = (0.3 - 0.9)$ TeV. Taking $T_{sp} = 80$ GeV, the bound becomes weaker: $M_{Z'} > (2.1 - 2.6)$ TeV.

Such a multi-TeV Z' is beyond the Tevatron II limit, but could be within the reach of LHC or future LC. Future collider sensitivities for Z' and also for N ($\tilde{N}$) have been analyzed in the literature. For LHC with the accumulated luminosity 100 fb^{-1}, a few hundred to a few $Z' \to NN$ events can be observed for $M_{Z'} = 3 - 5$ TeV. For LC at $E_{cm} = 3$ TeV with the luminosity 1000 fb^{-1}, 10^5 events will be obtained for $M_{Z'} = 3$ TeV [16]. With such large number of events, we will be able to study the properties of N and $\tilde{N}$. To cover more parameter space of our cosmological interest, an up-graded version of LHC or a higher energy option for LC will certainly be useful.

Apart from the well-studied Z' signals and dilepton (plus diHiggs) final states from $Z' \to NN$ [5], sneutrino-antisneutrino mixing and CP violating phenomena can be looked for to test the soft leptogenesis mechanism. Near the resonance point, $B \sim \Gamma$, order-one $\tilde{N}$-$\tilde{N}^\dagger$ oscillation effects are expected to occur in colliders, favoring the parameter space of $M_{Z'} \lesssim 3$ TeV and $\epsilon \sim 1$ which also requires a large CP phase, $\text{Im}(A)$. Following the decay $Z' \to \tilde{N}\tilde{N}^\dagger$,

the final states with a same-sign lepton (slepton) pair and chargino (charged Higgs) can occur through oscillations:

$$\tilde{N}\tilde{N}^{\dagger} \Rightarrow \tilde{N}\tilde{N}, \ \tilde{N}^{\dagger}\tilde{N}^{\dagger}$$
$$\rightarrow l^{\pm}l^{\pm}\chi^{\mp}\chi^{\mp} \tag{11}$$
$$\rightarrow \tilde{l}^{\pm}\tilde{l}^{\pm}H^{\mp}H^{\mp}.$$

As in the B-$\bar{B}$ mixing [17], let us define the oscillation parameters for sneutrinos: $x \equiv \Delta M_{\tilde{N}}/\Gamma = |B|M/\Gamma M_{\tilde{N}}$ and $y \equiv \Delta\Gamma/2\Gamma$ where $\Gamma = (\Gamma_1 + \Gamma_2)/2$ and $\Delta\Gamma = \Gamma_1 - \Gamma_2$, which can be measures in the mixing processes as we discuss. Various observables involving the oscillation parameters and the CP phase $\text{Im}(A)$ are as follows.

(i) Oscillation from the same-sign dileptons:

$$\frac{N(l^-l^-) + N(l^+l^+)}{N(l^+l^-) + N(l^-l^+)} = \frac{2r}{1+r^2}$$

where $r \equiv (x^2 + y^2)/(2 + x^2 - y^2)$.

(ii) Associated CP asymmetry:

$$\frac{N(l^-l^-) - N(l^+l^+)}{N(l^-l^-) + N(l^+l^+)} = \text{Im}\left(\frac{h^2 A}{4\pi B}\right).$$

(iii) Oscillation from the same-sign disleptons:

$$\frac{N(\tilde{l}^-\tilde{l}^-) + N(\tilde{l}^+\tilde{l}^+)}{N(\tilde{l}^-\tilde{l}^+) + N(\tilde{l}^+\tilde{l}^-)} = \frac{(1+r^2)a + r(1+a^2)}{(1+r^2)(1+a^2) + 2ra}$$

where $a \equiv |A/M|^2$.

(iv) The ratio of CP asymmetries:

$$\frac{N(\tilde{l}^-\tilde{l}^-) - N(\tilde{l}^+\tilde{l}^+)}{N(l^-l^-) - N(l^+l^+)} = \frac{|A|^4 - |M|^4}{M_{\tilde{N}}^4}.$$

If one measures the above quantities together with the decay rates and the sneutrino mass, one can basically reconstruct the amount of the cosmological lepton asymmetry given in Eq. (5) and obtain the final baryon asymmetry from Eq. (10), which could provide a way to study cosmology in colliders.

In conclusion, it is shown that, in the low energy $U(1)'$ models endowed with the seesaw mechanism, a successful leptogenesis can arise from supersymmetry breaking sector realizing the resonance behavior, $B \sim \Gamma$. Such a cosmological consideration puts a lower bound on the

Z' mass typically at multi TeV, assuming the standard gauge coupling strength. The bound is sensitive to the right-handed (s)neutrino mass as well as the spharelon rate which has a certain theoretical uncertainty. Further understanding on the spharelon dynamics would be required in the future confronting collider experiments. It is also pointed out that the soft leptogenesis mechanism of the model can be readily tested by the future observation of not only the Z' at multi-TeV range but also sneutrino-antisneutrino mixing and CP violation leading to distinct signatures of the same-sign dilepton (dislepton) and dichargino (diHiggs) production.

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